

Meson Correlation Function and Screening Mass in Thermal QCD

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Abstract

Analytical results for the spatial dependence of the correlation functions for all meson excitations in perturbative Quantum Chromodynamics, the lowest order, are calculated. The meson screening mass is obtained as a large distance limit of the correlation function. Our analysis leads to a better understanding of the excitations of Quark Gluon Plasma at sufficiently large temperatures and may be of relevance for future numerical calculations with fully interacting Quantum Chromodynamics.

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I. INTRODUCTION

In this work we study mesonic correlation functions in high temperature Quantum Chromodynamics (QCD) in the lowest order of the perturbative expansion. We follow the method developed by [1] where the simplest pseudoscalar case was calculated. Calculations are performed analytically with a special attention to the divergent parts, which need to be regularized and we adopt a Pauli-Villars [2] regularization scheme. The range of the correlation function is determined by the screening mass, which can be defined as the inverse screening length characterizing the exponential fall-off of the mesonic spatial correlator. For zero quark mass at the high temperature T the meson and baryon screening masses approach their ideal gas value, $2\pi T$ and $3\pi T$ [3, 4] respectively.

The pseudoscalar meson screening mass was calculated [1] in the noninteracting case $m_{scr} = 2\sqrt{\pi^2 T^2 + m^2}$, for finite quark mass m . Calculations in the interacting Quark Gluon Plasma (QGP) were done in nonrelativistic QCD [5–7] for all meson channels and in the Hard Thermal Loop Approximation (HTL) of the QCD [8] for the pseudo scalar one. The last, numerical approach, requires the analytical form of the mesonic spatial correlation functions in the free case in order to calculate the divergent integrals.

The main scope of this work is to provide the analytical expressions for the correlation functions for all meson channels to allow for future calculations of the screening masses in the interacting QGP.

The importance of such evaluation in order to identify the relevant degrees of freedom in hot and dense QCD was first pointed out in [9–11], where the screening masses of mesons and nucleons were calculated on the lattice. More recent results can be found in [12–23]. Though the evaluation of hadronic screening masses is a major achievement of lattice studies of the degrees of freedom characterizing the hot QCD (the large number of lattice sites available along the spatial directions allows one to study the large distance behaviour of the correlators), analytical approaches are not so common in the literature.

II. MESONIC SPATIAL CORRELATION FUNCTIONS

The lowest order quark loop contribution to the correlation function is defined by the expression

$$G_M(q) = -2N_c T \sum_n \int \frac{d^3p}{(2\pi)^3} \text{Tr} \left[\Gamma_M \frac{i}{\not{p} + \not{q} - m} \Gamma_M \frac{i}{\not{p} - m} \right] \quad (1)$$

from which we calculate spatial dependence for static plane-like perturbation

$$G_M(z) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} dq_z G_M(0, q_z^2) e^{iq_z z}. \quad (2)$$

Here M=S, PS, V, PV corresponds to the scalar, pseudoscalar, vector and pseudovector channel ($\Gamma_S = 1$, $\Gamma_{PS} = \gamma_5$, $\Gamma_V = \gamma_\mu$, $\Gamma_{PV} = \gamma_5 \gamma_\mu$), $q^\mu = (i\omega_m, \mathbf{q})$ is the external momentum, $\omega_m = 2m\pi iT$ is a boson Matsubara frequency. Summation runs over fermion Matsubara frequencies of quark internal momentum $p^\mu = (i\omega_n, \mathbf{p})$, $\omega_n = (2n+1)\pi iT$ and Tr denotes the trace over the spinor indices.

$$\begin{aligned} \text{Tr} \left[\Gamma_M \frac{i}{\not{p} + \not{q} - m} \Gamma_M \frac{i}{\not{p} - m} \right] &= \\ (Scalar) &= -\frac{4(m^2 + pq + p^2)}{(p^2 - m^2)((p+q)^2 - m^2)}, \\ (Pseudoscalar) &= -\frac{4(m^2 - pq - p^2)}{(p^2 - m^2)((p+q)^2 - m^2)}, \\ (Vector) &= -\frac{4(4m^2 - 2pq - 2p^2)}{(p^2 - m^2)((p+q)^2 - m^2)}, \\ (Pseudovector) &= -\frac{4(-4m^2 - 2pq - 2p^2)}{(p^2 - m^2)((p+q)^2 - m^2)}. \end{aligned} \quad (3)$$

We convert the sum over the frequencies [24] to two integrals. The first one, independent of T which is called a vacuum part (divergent one) and the second, T -dependent, called the matter part.

$$G_M = G_M^{vac} + G_M^{mat}. \quad (4)$$

Let's start from the vacuum part first

$$\begin{aligned}
G_S^{vac}(q) &= -8iN_c \int_{-i\infty}^{+i\infty} \frac{dp^0}{2\pi} \int \frac{d^3p}{(2\pi)^3} \frac{p^2 - q^2/4 + m^2}{[(p - q/2)^2 - m^2][(p + q/2)^2 - m^2]}, \\
G_{PS}^{vac}(q) &= 8iN_c \int_{-i\infty}^{+i\infty} \frac{dp^0}{2\pi} \int \frac{d^3p}{(2\pi)^3} \frac{p^2 - q^2/4 - m^2}{[(p - q/2)^2 - m^2][(p + q/2)^2 - m^2]}, \\
G_V^{vac}(q) &= 16iN_c \int_{-i\infty}^{+i\infty} \frac{dp^0}{2\pi} \int \frac{d^3p}{(2\pi)^3} \frac{p^2 - q^2/4 - 2m^2}{[(p - q/2)^2 - m^2][(p + q/2)^2 - m^2]}, \\
G_{PV}^{vac}(q) &= 16iN_c \int_{-i\infty}^{+i\infty} \frac{dp^0}{2\pi} \int \frac{d^3p}{(2\pi)^3} \frac{p^2 - q^2/4 + 2m^2}{[(p - q/2)^2 - m^2][(p + q/2)^2 - m^2]}. \tag{5}
\end{aligned}$$

After a Wick rotation ($p^0 \rightarrow ip_4$ and $q^0 \rightarrow iq_4$) in the Euclidean metric ($q_E^2 = q_4^2 + \mathbf{q}^2 = -q^2$) we can express (5) with two integrals

$$I_1(m) = 8N_c \int \frac{d^4p_E}{(2\pi)^4} \frac{1}{p_E^2 + m^2}, \tag{6}$$

$$I_2(m, q_E) = -4N_c \int \frac{d^4p_E}{(2\pi)^4} \frac{1}{[(p_E - q_E/2)^2 + m^2][(p_E + q_E/2)^2 + m^2]}, \tag{7}$$

where $d^4p_E = d^3p dp_4$. Both integrals are divergent and must be regularized. Following [1] we use Pauli-Villars regularization and define regularized functions:

$$I_1^R = \sum_{i=0}^N a_i I_1(M_i) \quad \text{and} \quad I_2^R(q_E) = \sum_{i=0}^N a_i I_2(M_i, q_E), \tag{8}$$

where $a_0 = 1$, $M_0 = m$ and higher masses are large. Coefficients a_i for $i > 0$ fulfill the following equations

$$\sum_{i=0}^N a_i = 0, \quad \sum_{i=0}^N a_i M_i^2 = 0, \quad \dots \quad \sum_{i=0}^N a_i M_i^{2(N-1)} = 0. \tag{9}$$

After some calculations we have

$$\begin{aligned}
I_1^R &= \frac{N_c}{2\pi^2} \sum_{i=0}^N a_i M_i^2 \ln M_i^2, \\
I_2^R(q_E) &= \frac{N_c}{2\pi^2} \sum_{i=0}^N a_i \left[\frac{2M_i}{q_E} \sqrt{1 + \left(\frac{q_E}{2M_i}\right)^2} \ln \left(\sqrt{1 + \left(\frac{q_E}{2M_i}\right)^2} + \frac{q_E}{2M_i} \right) + \ln M_i \right]. \tag{10}
\end{aligned}$$

The vacuum part of the correlation function can be written

$$\begin{aligned}
G_S^{vac}(q) &= -I_1^R - q_E^2 I_2^R(q_E) - 4m^2 I_2^R(q_E), \\
G_{PS}^{vac}(q) &= I_1^R + q_E^2 I_2^R(q_E), \\
G_V^{vac}(q) &= 2I_1^R + 2q_E^2 I_2^R(q_E) - 4m^2 I_2^R(q_E), \\
G_{PV}^{vac}(q) &= 2I_1^R + 2q_E^2 I_2^R(q_E) + 12m^2 I_2^R(q_E).
\end{aligned} \tag{11}$$

The next step is to obtain a z -dependent correlation function. Performing the Fourier transforms of (11) we find that in the complex q_E plane the function $I_2^R(q_E)$ has cuts for imaginary $q_E = ik$ from $\pm 2M_i$ to $\pm\infty$ respectively. Performing the integration (there is no contribution from I_1^R integral) the $q_E^2 I_2^R$ gives

$$\frac{N_c}{4\pi^2} \sum_{i=0}^N a_i \int_{2M_i}^{\infty} dk k \sqrt{k^2 - 4M_i^2} e^{-kz} = \frac{N_c}{\pi^2 z} \sum_{i=0}^N a_i M_i^2 K_2(2M_i z), \tag{12}$$

where K_2 is a modified Bessel function of a second kind. The I_2^R gives

$$\begin{aligned}
I_2^R(z) &= -\frac{N_c}{4\pi^2} \sum_{i=0}^N a_i \int_{2M_i}^{\infty} dk \frac{\sqrt{k^2 - 4M_i^2}}{k} e^{-kz} \\
&= -\frac{N_c}{8\pi^2} \sum_{i=0}^N a_i M_i G_{1,3}^{3,0} \left(z^2 M_i^2 \left| \begin{matrix} 1 \\ -\frac{1}{2}, 0, \frac{1}{2} \end{matrix} \right. \right),
\end{aligned} \tag{13}$$

where $G_{1,3}^{3,0}$ is a Meijer G function [25]. In both summations (12,13) the only first term survives and finally

$$\begin{aligned}
G_S^{vac}(z) &= -\frac{N_c}{\pi^2 z} m^2 K_2(2mz) + \frac{N_c}{2\pi^2} m^3 \mathcal{M}^{vac}(m, z), \\
G_{PS}^{vac}(z) &= \frac{N_c}{\pi^2 z} m^2 K_2(2mz), \\
G_V^{vac}(z) &= \frac{2N_c}{\pi^2 z} m^2 K_2(2mz) + \frac{N_c}{2\pi^2} m^3 \mathcal{M}^{vac}(m, z), \\
G_{PV}^{vac}(z) &= \frac{2N_c}{\pi^2 z} m^2 K_2(2mz) - \frac{3N_c}{2\pi^2} m^3 \mathcal{M}^{vac}(m, z),
\end{aligned} \tag{14}$$

where $\mathcal{M}^{vac}(m, z) = G_{1,3}^{3,0} \left(z^2 m^2 \left| \begin{matrix} 1 \\ -\frac{1}{2}, 0, \frac{1}{2} \end{matrix} \right. \right)$. The matter part of the correlation function is

$$G_M^{matt}(q) = -4iN_c \int_{-i\infty+\epsilon}^{+i\infty+\epsilon} \frac{dp^0}{2\pi} \int \frac{d^3 p}{(2\pi)^3} \frac{\text{Tr} \left[\Gamma_M \frac{i}{\not{p} + \not{q} - m} \Gamma_M \frac{i}{\not{p} - m} \right]}{e^{p^0/T} + 1}, \tag{15}$$

where Tr is evaluated in (3). After change of variables $p \rightarrow p' - q/2$, like in the vacuum part, and keeping the static limit ($q^0 = 0$) we evaluate the energy integral by deforming a contour of integration around the poles

$$p^0 = \omega_{\pm}(\mathbf{p}, \mathbf{q}) = \sqrt{m^2 + (\mathbf{p} \pm \mathbf{q}/2)^2} \quad (16)$$

and obtain

$$\begin{aligned} G_S^{matt}(q) &= -16N_c \int \frac{d^3p}{(2\pi)^3} \left[\frac{\mathbf{p}\mathbf{q} + q^2/2 + 2m^2}{4\omega_+ \mathbf{p}\mathbf{q}(e^{\omega_+/T} + 1)} + \frac{\mathbf{p}\mathbf{q} - q^2/2 - 2m^2}{4\omega_- \mathbf{p}\mathbf{q}(e^{\omega_-/T} + 1)} \right], \\ G_{PS}^{matt}(q) &= +16N_c \int \frac{d^3p}{(2\pi)^3} \left[\frac{\mathbf{p}\mathbf{q} + q^2/2}{4\omega_+ \mathbf{p}\mathbf{q}(e^{\omega_+/T} + 1)} + \frac{\mathbf{p}\mathbf{q} - q^2/2}{4\omega_- \mathbf{p}\mathbf{q}(e^{\omega_-/T} + 1)} \right], \\ G_V^{matt}(q) &= -32N_c \int \frac{d^3p}{(2\pi)^3} \left[\frac{\mathbf{p}\mathbf{q} + q^2/2 - m^2}{4\omega_+ \mathbf{p}\mathbf{q}(e^{\omega_+/T} + 1)} + \frac{\mathbf{p}\mathbf{q} - q^2/2 + m^2}{4\omega_- \mathbf{p}\mathbf{q}(e^{\omega_-/T} + 1)} \right], \\ G_{PV}^{matt}(q) &= -32N_c \int \frac{d^3p}{(2\pi)^3} \left[\frac{\mathbf{p}\mathbf{q} + q^2/2 + 3m^2}{4\omega_+ \mathbf{p}\mathbf{q}(e^{\omega_+/T} + 1)} + \frac{\mathbf{p}\mathbf{q} - q^2/2 - 3m^2}{4\omega_- \mathbf{p}\mathbf{q}(e^{\omega_-/T} + 1)} \right]. \end{aligned} \quad (17)$$

Then we simplify by changing the variables $\mathbf{p}' \rightarrow \mathbf{p} + \mathbf{q}/2$ in the first part of the sum and $\mathbf{p}' \rightarrow \mathbf{p} - \mathbf{q}/2$ in the second part and integrate over angles.

$$\begin{aligned} G_S^{matt}(q) &= \frac{N_c}{\pi^2} \int_0^\infty dp p^2 \frac{1}{\omega_p} \frac{1}{e^{\omega_p/T} + 1} \{A_1(p, q) + 2m^2 A_2(p, q)\}, \\ G_{PS}^{matt}(q) &= -\frac{N_c}{\pi^2} \int_0^\infty dp p^2 \frac{1}{\omega_p} \frac{1}{e^{\omega_p/T} + 1} \{A_1(p, q)\}, \\ G_V^{matt}(q) &= -\frac{2N_c}{\pi^2} \int_0^\infty dp p^2 \frac{1}{\omega_p} \frac{1}{e^{\omega_p/T} + 1} \{A_1(p, q) - m^2 A_2(p, q)\}, \\ G_{PS}^{matt}(q) &= -\frac{2N_c}{\pi^2} \int_0^\infty dp p^2 \frac{1}{\omega_p} \frac{1}{e^{\omega_p/T} + 1} \{A_1(p, q) + 3m^2 A_2(p, q)\}, \end{aligned} \quad (18)$$

where $\omega_p = \sqrt{p^2 + m^2}$, $p = |\mathbf{p}|$, $q = |\mathbf{q}|$ and

$$\begin{aligned} A_1(p, q) &= 4 + \frac{q}{p} \ln \left| \frac{2p - q}{2p + q} \right|, \\ A_2(p, q) &= \frac{2}{pq} \ln \left| \frac{2p - q}{2p + q} \right|. \end{aligned} \quad (19)$$

Since the whole q dependence is in the A_1 and A_2 functions, we only need the Fourier transform, see (2), of A_1 and A_2

$$\begin{aligned} A_1(p, z) &= \frac{1}{2\pi} \int_{-\infty}^{+\infty} dq \left\{ 4 + \frac{q}{p} \ln \left| \frac{2p-q}{2p+q} \right| \right\} e^{iqz}, \\ A_2(p, z) &= \frac{1}{2\pi} \int_{-\infty}^{+\infty} dq \frac{2}{pq} \ln \left| \frac{2p-q}{2p+q} \right| e^{iqz}. \end{aligned} \quad (20)$$

In order to perform the integration we use the identity

$$\ln \left| \frac{2p-q}{2p+q} \right| = \lim_{\epsilon \rightarrow 0} \frac{1}{2} \ln \frac{(2p-q)^2 + \epsilon^2}{(2p+q)^2 + \epsilon^2}. \quad (21)$$

The function (21) has two cuts on the upper half plane $q = ik + 2p$ and $q = ik - 2p$. Fourier transform of A_1 is easy to obtain

$$\begin{aligned} A_1(p, z) &= \lim_{\epsilon \rightarrow 0} \frac{1}{2\pi} \int_{-\infty}^{+\infty} dq \left\{ 4 + \frac{q}{2p} \ln \frac{(2p-q)^2 + \epsilon^2}{(2p+q)^2 + \epsilon^2} \right\} e^{iqz} \\ &= \frac{\sin(2pz)}{pz^2} - \frac{2 \cos(2pz)}{z}. \end{aligned} \quad (22)$$

The second case, A_2 , is more complicated and we finish with two integrals

$$\begin{aligned} A_2(p, z) &= \lim_{\epsilon \rightarrow 0} \frac{1}{2\pi} \int_{-\infty}^{+\infty} dq \frac{1}{pq} \ln \frac{(2p-q)^2 + \epsilon^2}{(2p+q)^2 + \epsilon^2} e^{iqz} \\ &= -\lim_{\epsilon \rightarrow 0} \frac{1}{p} \left[e^{2ipz} \int_{\epsilon}^{\infty} \frac{dk}{ik+2p} e^{-kz} - e^{-2ipz} \int_{\epsilon}^{\infty} \frac{dk}{ik-2p} e^{-kz} \right], \end{aligned} \quad (23)$$

where

$$\lim_{\epsilon \rightarrow 0} \int_{\epsilon}^{\infty} \frac{dk}{ik \pm 2p} e^{-kz} = -ie^{\mp 2ipz} \left[\Gamma(0, \mp 2ipz) + \ln \left(\frac{\pm i}{2p} \right) - \ln(z) + \ln(\mp ipz) \right] \quad (24)$$

so that finally we have

$$A_2(p, z) = \frac{i}{p} [\Gamma(0, -2ipz) - \Gamma(0, 2ipz)], \quad (25)$$

where $\Gamma(0, ix)$ is the incomplete gamma function. In order to obtain the z -dependent correlation function in the matter case we have to insert (22,23) into (18). To perform integration over p we can use the identity

$$\frac{1}{e^{\omega_p/T} + 1} = \frac{1}{2} - \sum_{l=-\infty}^{\infty} \frac{\omega_p T}{(2l+1)^2 \pi^2 T^2 + \omega_p^2}. \quad (26)$$

We then obtain

$$\begin{aligned}
\int_0^\infty \frac{p^2 dp}{\omega_p} \left[\frac{1}{2} - \sum_{l=-\infty}^\infty \frac{\omega_p T}{(2l+1)^2 \pi^2 T^2 + \omega_p^2} \right] A_1(p, z) &= \frac{m^2}{z} K_2(2mz) - \\
-\frac{\pi T}{2z^2} \sum_{l=-\infty}^\infty e^{-2z\sqrt{(2l+1)^2 \pi^2 T^2 + m^2}} \left(2z\sqrt{(2l+1)^2 \pi^2 T^2 + m^2} + 1 \right) &= \\
&= \frac{m^2}{z} K_2(2mz) - A_1^T(p, z), \tag{27}
\end{aligned}$$

for A_2 the first term of the expansion (26) gives

$$\int_0^\infty \frac{p^2 dp}{\omega_p} \frac{1}{2} A_2(p, z) = -\frac{1}{4\pi z} \Re \left[G_{2,4}^{4,1} \left(-z^2 M^2 \left| \begin{array}{c} \frac{1}{2}, \frac{3}{2} \\ 0, \frac{1}{2}, \frac{1}{2}, 1 \end{array} \right. \right) \right] = -\frac{1}{4\pi z} \mathcal{M}^{matt}(m, z). \tag{28}$$

For each element of the sum (26) (for every l and T) we have

$$\int_0^\infty \frac{T p^2 dp}{(2l+1)^2 \pi^2 T^2 + \omega_p^2} A_2(p, z) = 0. \tag{29}$$

Collecting (27,28,29) and (18) we have

$$\begin{aligned}
G_S^{matt}(z) &= \frac{N_c m^2}{\pi^2 z} K_2(2mz) - \frac{N_c}{\pi^2} A_1^T(p, z) - \frac{N_c m^2}{2\pi^3 z} \mathcal{M}^{matt}(m, z), \\
G_{PS}^{matt}(z) &= -\frac{N_c m^2}{\pi^2 z} K_2(2mz) + \frac{N_c}{\pi^2} A_1^T(p, z), \\
G_V^{matt}(z) &= -\frac{2N_c m^2}{\pi^2 z} K_2(2mz) + \frac{2N_c}{\pi^2} A_1^T(p, z) - \frac{N_c m^2}{2\pi^3 z} \mathcal{M}^{matt}(m, z), \\
G_{PS}^{matt}(z) &= -\frac{2N_c m^2}{\pi^2 z} K_2(2mz) + \frac{2N_c}{\pi^2} A_1^T(p, z) + \frac{3N_c m^2}{2\pi^3 z} \mathcal{M}^{matt}(m, z).
\end{aligned} \tag{30}$$

Taking into account the relation

$$\Re \left[G_{2,4}^{4,1} \left(-x^2 \left| \begin{array}{c} \frac{1}{2}, \frac{3}{2} \\ 0, \frac{1}{2}, \frac{1}{2}, 1 \end{array} \right. \right) \right] = \pi x G_{1,3}^{3,0} \left(x^2 \left| \begin{array}{c} 1 \\ -\frac{1}{2}, 0, \frac{1}{2} \end{array} \right. \right) \tag{31}$$

we have

$$\mathcal{M}^{matt}(m, z) = \pi m z \mathcal{M}^{vac}(m, z) \tag{32}$$

and finally after some cancellations

$$\begin{aligned}
G_S(z) &= \frac{-N_c T}{2\pi z^2} \sum_{l=-\infty}^{\infty} e^{-2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2}} \left(2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2} + 1 \right) \\
G_{PS}(z) &= \frac{N_c T}{2\pi z^2} \sum_{l=-\infty}^{\infty} e^{-2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2}} \left(2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2} + 1 \right) \\
G_V(z) &= \frac{N_c T}{\pi z^2} \sum_{l=-\infty}^{\infty} e^{-2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2}} \left(2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2} + 1 \right) \\
G_{VS}(z) &= \frac{N_c T}{\pi z^2} \sum_{l=-\infty}^{\infty} e^{-2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2}} \left(2z\sqrt{(2l+1)^2\pi^2 T^2 + m^2} + 1 \right)
\end{aligned} \tag{33}$$

The main contribution to the summations comes from two elements $l = 0$ and $l = -1$. In the $z \rightarrow \infty$ limit all channels give the same asymptotic

$$G(z) \xrightarrow{z \rightarrow \infty} \frac{\text{const}}{z} e^{-2\sqrt{\pi^2 T^2 + m^2} z} = \frac{\text{const}}{z} e^{-m_{scr} z} \tag{34}$$

where we identified the screening mass m_{scr}

$$m_{scr} = 2\sqrt{\pi^2 T^2 + m^2}, \tag{35}$$

which for massless quarks has a simple form $m_{scr} = 2\pi T$.

III. CONCLUSIONS

In this work we have shown that due to the specific cancellations of Meijer G functions the final formulas for the spatial correlation function are compact and differ only with the sign or some numerical constants between different channels in the weakly interacting QGP.

The screening mass is the quantity which governs the large-distance exponential decay of the correlations of mesonic current operators and is of course equal, in our system, for all mesons and has a value $m_{scr} = 2\sqrt{\pi^2 T^2 + m^2}$, where m is the free quark mass.

We have explicitly evaluated the analytical form of the spatial, z -dependent mesonic correlation functions for all mesonic channels, which is of great importance for future numerical calculations of the effective masses of mesons in the fully interacting QGP, especially that in the literature there is a discrepancy between lattice [12–23] and non-lattice [5, 6, 8] calculations of the mesonic screening masses. In the lattice calculations in the high temperature limit, screening masses approach the non interacting value from below, whereas in non lattice calculations they approach from above.

In papers [5, 6] the analytic predictions for the screening masses, related to various quark-antiquark excitations at high temperatures, are determined in the nonrelativistic 3-dimensional QCD effective theory with the next-to-leading-order perturbative corrections. In paper [8] the numerical calculation in the HTL approximation, in the pseudo scalar channel only, was performed with the help of a new technique of numerical regularization of divergent integrals with the usage of the analytical formulas of the spatial correlation functions from the non interacting case, which we provide in this work.

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